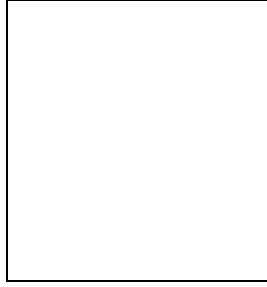


DIFFUSIVE SHOCK ACCELERATION: THE FERMI MECHANISM

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The mechanism of diffusive Fermi acceleration at collisionless plasma shock waves is widely invoked in astrophysics to explain the appearance of non-thermal particle populations in a variety of environments, including sites of cosmic ray production, and is observed to operate at several sites in the heliosphere. This review outlines the principal results from the theory of diffusive shock acceleration, focusing first on how it produces power-law distributions in test-particle regimes, where the shock dynamics are dominated by the thermal populations that provide the seed particles for the acceleration process. Then the importance of non-linear modifications to the shock hydrodynamics by the accelerated particles is addressed, emphasizing how these subsequently influence non-thermal spectral formation.

1 Introduction

The concept of diffusive acceleration of particles in space plasmas has been around for nearly five decades. Fermi¹ (1949) first postulated that cosmic rays could be produced via diffusion between collisions between interstellar clouds. If such clouds have predominantly random directions of motion, then the frequency of “head-on” collisions between the cosmic rays and the clouds would exceed the rate of “tail” encounters, leading to a net acceleration that is diffusive in energy. The elegance of Fermi’s idea was founded on the fact that, when particles are confined, such a diffusive process naturally produces power-law¹ cosmic ray distributions, thereby modelling the observations well. It was subsequently realized that shocks in space plasmas could also provide such diffusive acceleration in an efficient manner, tapping the dissipative potential of the flow discontinuity by transferring the shock’s kinetic energy to non-thermal populations both upstream and downstream of the shock, at the same time as heating the downstream gas. Thus the notion of diffusive shock acceleration, or the *first-order Fermi* mechanism, was born.

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Fermi acceleration at shocks is applied to many physical systems in the solar system, in our galaxy, and throughout the universe, primarily because (i) of its great efficiency for generating non-thermal ions and electrons, and (ii) since inferences of non-thermal particles abound in astrophysics in addition to the direct observations of such populations in the heliosphere. The closest environment for studying such acceleration is the Earth's bow shock, a standing shock formed by the Earth's passage through the solar wind. Observations of non-thermal ions on either side of quasi-parallel portions (where the magnetic field is normal to the shock interface) of the bow shock (e.g. Gosling et al.²; Ellison et al.³) are a testament to the efficiency of diffusive shock acceleration. Bow shocks of other planets and comets also exhibit particle acceleration. Throughout the solar system, a whole host of travelling *interplanetary* shocks and forward/reverse shock pairs in *co-rotating interaction regions* propagate through space, all providing unequivocal evidence for acceleration of protons, alpha particles and other ions (e.g. Gosling et al.⁴; Baring et al.⁵). These non-linear disturbances exemplify the complexity of the plasma and magnetic structure of the solar wind in its journey to the outer solar system, at which point it is commonly believed to encounter at some distance greater than 70 AU a solar wind *termination shock*: this defines the boundary between the heliosphere and the surrounding interstellar medium. Evidence for the existence of such a structure is still circumstantial.

In astrophysics, the number of sites for shock acceleration burgeons. All these are based on inferences obtained from non-thermal radiation, given that in situ measurements are impossible. Stellar winds (similar to the solar wind), supernova remnants (SNRs), accreting X-ray binaries and pulsar wind termination shocks are some examples of sources of shock acceleration within the Milky Way. For example, supernova remnants commonly exhibit large "surfaces" of sharp intensity variations in radio and X-ray images (e.g. see references in the catalogue of Green⁶). Such interfaces, with a variety of spatial morphologies, are naturally expected as the supernova ejecta in the remnant ploughs into the surrounding interstellar medium (ISM). Concurrent with the existence of such interfaces is non-thermal radio synchrotron radiation in remnants, and the recent discovery of non-thermal X-rays from remnants (e.g. Koyama et al.⁷). SNRs have long been thought to be the preferred site for galactic cosmic ray generation for energies up to the so-called *knee* in the spectrum at $\sim 10^{15}$ eV (e.g. see Lagage and Cesarsky⁸), where significant deviations from power-law behaviour are observed. The highest energy cosmic rays ($\gtrsim 10^{18}$ eV) are believed to be of extragalactic origin, where the large size scales permit acceleration to energies beyond the capabilities of galactic environments. They could be provided by shocks associated with jets and hot spots from radio galaxies, whose jets display a clumpiness (knots) reminiscent of shocks (e.g. Blandford and Eichler⁹). There are also proponents of the idea that the most energetic cosmic rays at $\sim 10^{20}$ eV may be caused by shocks in gamma-ray bursts (Waxman¹⁰; Milgrom and Usov¹¹). The prevalence of non-thermal X-ray emission in Seyfert galaxies and gamma-rays from blazars promotes the notion that the central regions of active galaxies possess shocks that efficiently accelerate electrons and/or ions. The underlying theme that should be borne in mind is that wherever shocks are seen, there is evidence of non-thermal particles: this strongly suggests that shock acceleration is ubiquitous in the universe.

The modern era of shock acceleration theory began with a collection of seminal papers in 1977-78 (Krymsky¹²; Bell¹³; Axford, Lear and Skadron¹⁴; Blandford and Ostriker¹⁵) that outlined the basic properties of the process. Since then the field has grown substantially, with numerous approaches being developed, each with their different attributes and limitations. In this review, the basic properties of shock acceleration theory that are *most relevant* to astrophysical modelling and data comparison will be reviewed, starting with a brief outline of the principal of the Fermi mechanism. Then the main result of the linear or *test particle* theory of acceleration, namely the emergence of canonical power-law distributions, will be derived, and a discussion of simulation results and predictions of acceleration efficiencies will be reviewed. The importance of diffusion of particles perpendicular to the ambient field for quasi-perpendicular shocks will

be discussed. Finally the importance of non-linear feedback between the acceleration process and the shock hydrodynamics will be addressed, an effect that is crucial to the understanding of many astrophysical shocks. There are many discussions of shock acceleration in the literature, however three principal reviews with complementary focuses and approaches come to mind. Drury¹⁶ provides an in-depth analytic description of linear and non-linear shocks, tailored for the specialist, while Blandford and Eichler⁹ adopt a slightly tempered analytic approach and connect closely with astrophysical observations. Jones and Ellison¹⁷ emphasize the more recent model developments and data/theory comparisons afforded by computer simulations.

2 The Principal of the Fermi Mechanism

The Fermi acceleration mechanism is always applied to so-called *collisionless* shocks, i.e. those non-linear disturbances that have energy and momentum transfer between particles mediated purely by plasma processes, with Coulomb scattering being negligible. Such conditions arise in most astrophysical environments. It is instructive to review the principal of how the Fermi mechanism operates. Consider a flow defined by speeds u_1 and u_2 ($< u_1$) on different sides of the shock (see Fig. 1). Suppose particles of speed v_0 (in the rest frame of the shock) start off on the upstream (u_1) side of the shock, and diffuse around via “collisions” with plasma magnetic turbulence until they eventually cross the shock and move around downstream (u_2). Such diffusion tends to isotropize the angular distribution of the particles in the frame in which the upstream plasma is at rest. After a period in which the particles have experienced a plasma of speed u_1 , the particles now collide with magnetic turbulence that is associated with, or generated by, the downstream plasma. If this plasma in turn tends to isotropize the particles elastically, then this test population effectively “sees” a plasma moving *towards* it (with speed $\sim |u_1 - u_2|$) upon arrival downstream. Elementary kinematics leads to the result that the process of quasi-isotropization then yields a net *increase* in the average particle speed in the rest frame of the shock interface. This kinematic guarantee of an increase in energy is akin to the gain that a photon experiences in inverse Compton scattering collisions with relativistic electrons, an effect that relies on photon quasi-isotropization in the electron rest frame.

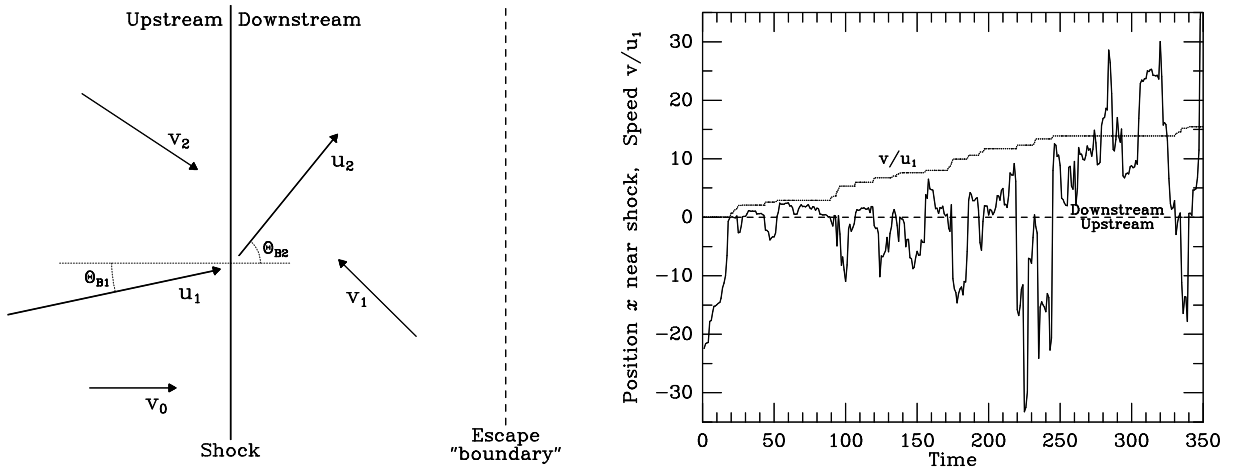


Figure 1: A schematic depiction (left) of particle motion in the environs of a shock. The plasma flow speeds are u_1 (upstream) and u_2 (downstream), and the mean accelerated particle speeds v_i after i shock crossings are ordered according to $v_0 < v_1 < v_2$, etc. The dotted line denotes a downstream “escape boundary” that is relevant to the discussion on universal power-laws in Section 3 below. On the right is output from a Monte Carlo simulation of first-order Fermi acceleration (Baring et al.¹⁸), illustrating how the particle speed v increases monotonically with time as many shock crossings are encountered, until it escapes the shock (dotted line). Notice that the diffusion is on larger scales for higher particle speeds.

Eventually, some of the particles will return to the upstream side of the shock, and these will see the upstream plasma moving, on average, towards them. Hence the process is repeated, and diffusion that leads to significant isotropization in the local plasma frame will yield particle speeds that are on average higher than those previously obtained on the downstream side of the shock. Sequential diffusion back and forth across the shock always leads to increases in particle speed, so that many shock crossing cycles afford significant acceleration. This is the principal of the *first-order* Fermi mechanism, where energy gains are always positive, and spatial diffusion allows the shock to dissipate its energy via its associated turbulence to a non-thermal particle population. This situation is depicted in Fig. 1, both schematically, and also in an example where an actual Monte Carlo simulation¹⁸ was used to demonstrate the coupling between shock crossings and energy increase. The velocity increase, for non-relativistic particles, is proportional to $\sim |u_1 - u_2|$, a feature that is evident in Fig. 1, so that many cycles are required for particles to achieve high energies. For a simple visualization, the Fermi mechanism can be likened to specular reflection between two converging mirrors, for example a ball bouncing elastically back and forth between two walls that approach each other.

Before moving to a discussion of aspects of test-particle acceleration, a brief summary of approaches to shock acceleration theory is desirable. Early in the 1980s, *two-fluid* models (see Drury’s¹⁶ review) were developed, treating the cosmic rays and thermal gas as separate entities, and exploring the hydrodynamics of shocked flows subject to non-linear shock acceleration. These are extremely useful for time-dependent applications, but contain little spectral information. Differential equations that describe the convective and diffusive parts of particle motion have been used in a variety of situations with a variety of assumptions; they handle time-dependence and spectral information well, but have difficulty injecting particles from the thermal population. Plasma codes¹⁷ compute particle motions and field structure using the Newton-Lorentz and Maxwell’s equations, and hence are physically the most detailed approach to acceleration theory. They come in full and hybrid (where electrons are treated as a background fluid) varieties, are great for studying wave properties, but have a severely CPU-limited spectral dynamic range. Finally, kinetic Monte Carlo simulations¹⁷, of which several examples are presented here, describe the convective and diffusive particle motions using a prescribed scattering or diffusion law. They self-consistently inject particles from thermal energies, have excellent spectral information and dynamic range, and hence are ideal for comparison with experimental data. The Monte Carlo technique is presently limited to time-independent situations.

3 Acceleration in Test-Particle Regimes

The ions accelerated by the Fermi mechanism are called test particles when they do not have sufficient pressure (i.e. energy density) to influence the hydrodynamics of the shocked plasma, and hence “go along for the ride.” The description of the shock acceleration mechanism is then linear and therefore comparatively simple. Moreover, when the test particle speeds far exceed that of the shock, or equivalently u_1 and u_2 , the system possesses no velocity scale (not true of non-linear acceleration) so that the natural solution for the distribution function of the non-thermal particles (also frequently referred to as the cosmic ray component) is a power-law distribution. This is the foremost property of the test-particle, or linear, acceleration regime, and it is instructive to derive the “canonical” power-law index.

Let the distribution function for particles of speed v and momentum $p = mv$ in some frame of reference (say the shock frame) be $f(p)$, and the cumulative distribution between p and infinity be $\mathcal{F}(p)$. Assume that $v \gg u_1$ and that the particles are isotropic; under these conditions, the choice of reference frame among the upstream and downstream plasma frames or the shock frame is immaterial—isotropy is attained in all three. Then it is possible to write down

the differential equation governing $\mathcal{F}(p)$: in a steady-state scenario, the system is described by

$$0 \equiv t_{cyc} \frac{d\mathcal{F}}{dt} = -\langle \Delta p \rangle \frac{\partial \mathcal{F}}{\partial p} - P_{esc} \mathcal{F} \quad , \quad \mathcal{F}(p) = \int_p^\infty f(p_1) dp_1 \quad . \quad (1)$$

Here t_{cyc} is the mean time for a cycle of *two* transmissions through the shock, e.g. upstream to downstream and then back upstream again, and $\langle \Delta p \rangle$ is the mean (positive) net momentum gain in a cycle. Hence the $-\langle \Delta p \rangle \partial \mathcal{F} / \partial p$ term represents the supply of particles to higher energies via the acceleration process, and involves a derivative with respect to p because the momentum gains are approximately differential: $\langle \Delta p \rangle \sim m|u_1 - u_2| \ll p$. The negative sign appears in this term due to the use of a cumulative distribution. The remaining term represents a loss of particles from each cycle (with probability P_{esc}) beyond some downstream reference boundary that is depicted in Fig. 1. The existence of such a loss is a consequence of the net sense of the plasma flow in the downstream direction; when $v \gg u_1$ then the spatial diffusion of the particles back and forth across this boundary dominates convection from the flow and the probability of escape is small ($P_{esc} \ll 1$). This escape term is proportional to $\mathcal{F}(v)$ since particles that are lost at a particular speed cannot contribute further to the acceleration process.

The form of Eq. (1) corresponds to the first-order Fermi process, so named because the acceleration is first-order in velocity differentials with the equation having only friction terms in p . Second-order Fermi acceleration, for which there is diffusion in p , can also occur and yield power-law distributions. It is generally a minor contribution unless the shock speed is almost as low as the Alfvén speed $v_A = B/\sqrt{4\pi\rho}$, which defines the “speed” of the magnetic disturbances in the plasma that effect particle scattering and diffusion. Hence Alfvén speeds of the order of u_1 lead to extra terms in the Fokker-Planck expansion that has its friction term appearing in Eq. (1). Solutions to this equation are easily obtained. For isotropy of particles, the probability of downstream escape of particles with $v \gg u_1$ from a cycle is $P_{esc} \approx 4u_1/v$ for *plane-parallel* shocks where the field is normal to the discontinuity interface. The average momentum gain in a cycle for isotropic populations is $\langle \Delta p \rangle \approx 4u_1(u_1 - u_2)p/(3u_2v)$. Hence, defining the shock *compression ratio* as $r = u_1/u_2$, Eq. (1) admits power-law solutions (for $u_1 \ll c$):

$$f(v) \propto p^{-\sigma} \quad , \quad \sigma = 1 + \frac{pP_{esc}}{\langle \Delta p \rangle} = \frac{r+2}{r-1} \quad , \quad r = \frac{u_1}{u_2} \quad . \quad (2)$$

This is the canonical or universal power-law distribution that is frequently invoked in shock acceleration applications to astrophysics, and is valid both for non-relativistic and relativistic particles. The power-law index σ depends only on the compression ratio r , an elegant feature of the first-order Fermi mechanism, and non-relativistic plasma shocks (of adiabatic index $\gamma = 5/3$) with large sonic Mach numbers (so-called *strong* shocks) have $r = 4$ and therefore $\sigma = 2$. If the plasma speeds $u_{1,2}$ used in computing the compression ratio are interpreted as flow velocity components along the shock normal, then the result in Eq. (2) extends to *oblique* shocks (e.g. Drury¹⁶), where the mean magnetic field \mathbf{B} makes significant angles with the shock normal (see Fig. 1). Furthermore, the canonical power-law is applicable regardless of the relative contributions of particle diffusion along and orthogonal to \mathbf{B} (Jones¹⁹). Alternative derivations of the universal power-law are presented by Bell³, Drury¹⁶, Blandford and Eichler⁹ and Jones and Ellison¹⁷. Note that while second-order Fermi acceleration also yields power-laws, there is no coupling between σ and the compression ratio r , primarily because σ depends on two independent parameters: the speed of the scatterers (i.e. magnetic turbulence) in the flow frame, and the residence time in the acceleration region. Hence there is **no universality of the power-law** for 2nd-order acceleration, underlining the beauty of the first-order process.

In the shock acceleration mechanism, the canonical power-laws are attained only at high energies, so that a more complex (i.e. non-analytic) description becomes necessary when $v \sim u_1$. Measurements of shock-associated non-thermal ions in the heliosphere clearly indicate that low

energy non-thermal ions are drawn directly from the thermal component of the plasma. In particular, it seems that all ions are subject to the same diffusive behaviour, with no peculiar situation arising at either low or high energies. Hence, a coherent treatment of Fermi acceleration must give similar weight to thermal and superthermal particles; this is a nice feature of some of the computer simulations reviewed by Jones and Ellison¹⁷. Such simulations provide insight into the relative abundances of thermal and non-thermal ion components; these reflect the inherent efficiency of the acceleration process. Therefore, we turn our attention now to survey results from simulations that address salient properties of linear shock acceleration theory, such as the dependence of the shape of the distribution and the acceleration efficiency on shock obliquity, particular plasma characteristics and the description of scattering.

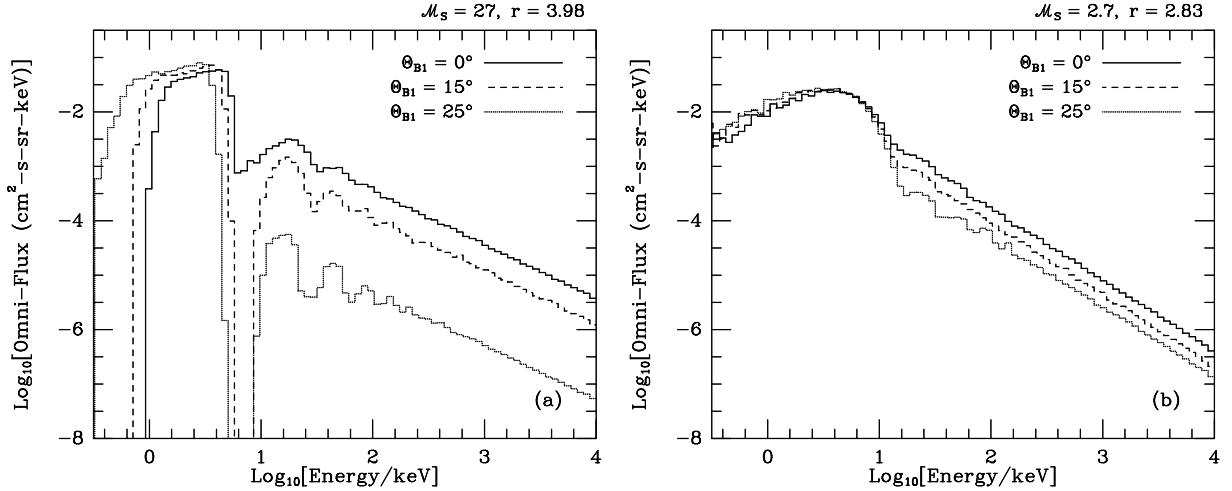


Figure 2: Omni-directional flux distributions for different shock obliquities, normalized to unity and measured downstream of (a) high and (b) low sonic Mach number shocks of speed 1000 km/sec, as obtained^{18,23} from kinematic Monte Carlo simulations of Fermi acceleration. Particles were injected into the upstream region, with temperatures $T_1 = 10^5$ K ($\mathcal{M}_s = 27$) and $T_1 = 10^7$ K ($\mathcal{M}_s = 2.7$), and were allowed to escape from the shock either far downstream or when their energies exceeded 10^4 keV. A dramatic reduction in injection efficiency with increasing obliquity is evident in the high Mach number case (a) with the normalization dropping by about a factor of 100 between $\Theta_{B1} = 0^\circ$ and 25° . The power-law indices closely approximate those in Eq. (2).

Although the power-law index of the accelerated particles is independent of shock obliquity, the acceleration efficiency is *strongly* dependent on the angle Θ_{B1} that the upstream magnetic field makes to the shock normal. This is illustrated in Fig. 2, where “spectral” results from a kinematic Monte Carlo approach to shock acceleration (Baring et al.¹⁸) are shown. The distributions, obtained in a region just downstream of the shock, exhibit prominent thermal components at low energies, and canonical power-law tails at high energies, with structure in between (for $u_1 \lesssim v \lesssim 10u_1$). The prominence of this structure originates in the simplistic nature of the particle trajectories assumed in these calculations, and is largely smeared out by more detailed approaches (e.g. Ellison et al.²⁰) that include particle gyromotions. Notwithstanding, the structure signifies successive shock crossings, ordered according to increasing energy, and merges smoothly into the power-law tail when $v \gg u_1$. In the left hand panel, where the upstream plasma is very cold (and the sonic Mach number $\mathcal{M}_s \approx u_1 \sqrt{m_p/(\gamma k T_1)}$ is large) the thermal peaks are not truly Maxwellian, but rather are defined by the kinematics of the simulation: a more complicated picture of shock dissipation and heating would yield almost Maxwellian shapes (e.g. see results from detailed plasma simulations in Giacalone et al.²¹, Scholer et al.²²). In the right hand panel, the higher temperature (lower \mathcal{M}_s) smears out the structure. Note that for lower \mathcal{M}_s , the hotter upstream plasma produces weaker shocks (lower r) and therefore steeper distributions. The normalization of the non-thermal component relative to the thermal population describes the overall efficiency of acceleration. This is seen to drop dramatically for

large Mach numbers when significant obliquities Θ_{B1} are realized. The reason for this (Baring et al.^{18,23}) is that as the obliquity increases, the mean increment in the particle speed $\langle \Delta v \rangle$ upon first crossing of the shock declines, and turns out to approach zero in the limit of *quasi-perpendicular* (i.e. $\Theta_{B1} \approx 90^\circ$) shocks. When the component of this increment along the mean field direction drops below u_2 , particles cannot return to the upstream side of the shock so that injection into the Fermi process fails. This occurs²³ when $\Theta_{B1} \gtrsim 30^\circ$ for cold upstream plasmas and at $\Theta_{B1} \gtrsim 55^\circ$ for the warm plasma case in right hand panel of Fig. 2.

This termination of injection poses a serious problem for highly oblique and quasi-perpendicular shocks. These disturbances abound in the heliosphere as travelling interplanetary shocks and discontinuities associated with co-rotating interaction (CIR) regions. Due to the tightness of the winding of the spiral heliospheric field in the solar wind, most interplanetary and CIR shocks are highly oblique. Yet they are observed to be prolific accelerators of ions (for recent Ulysses observations see, e.g., Ogilvie, et al.²⁴; Gloeckler, et al.²⁵; Baring et al.⁵). Therefore, the theory leading to the distributions in Figure 2 lacks an ingredient that is key to the acceleration process in highly oblique shocks — this ingredient is *cross-field diffusion*. The simulation results in Figure 2 included only a description of particle diffusion along the local magnetic field lines. Clearly turbulent plasmas provide diffusion across the mean direction of \mathbf{B} , as is deduced from spacecraft observations of interplanetary magnetic field power spectra (e.g. Forman et al.²⁶; Valdes-Garcia et al.²⁷). To add this to Monte Carlo simulations effectively requires specification of a *spatial diffusion coefficient* perpendicular to the mean local (i.e. upstream or downstream) field, κ_\perp . Using the standard kinetic theory result for field turbulence (e.g. Axford²⁸; Forman et al.²⁶), this can be related to the diffusion coefficient along the field, κ_\parallel , via

$$\kappa_\perp = \frac{\kappa_\parallel}{1 + (\lambda/r_g)^2} \quad . \quad (3)$$

Here λ is the mean free path for particle scattering (i.e. diffusion) along the field, which is related to the spatial diffusion coefficient by $\kappa_\parallel = \lambda v/3$. Also, $r_g = pc/(qB)$ is the gyroradius of a particle of momentum p and charge q . In the limit of small mean free paths, $\lambda \sim r_g$, the diffusion satisfies $\kappa_\perp \sim \kappa_\parallel$ and is quasi-isotropic. This defines the *Bohm* diffusion limit where the field turbulence is strong and no direction is preferred for particle diffusion; in this limit the Fermi acceleration of ions should be more or less equivalent for quasi-parallel ($\Theta_{B1} \sim 0^\circ$) and quasi-perpendicular ($\Theta_{B1} \sim 90^\circ$) shocks. Values of $\lambda \lesssim r_g$ are physically unrealistic.

The inclusion of cross-field diffusion short-circuits the termination of injection, provided that it can return to the shock particles that have entered the downstream region for the first time. Since particles diffuse of the order of their gyroradius r_g perpendicular to the field every time they are “scattered” along the field (i.e. on the scale of λ), simple geometry considerations imply that the criterion for the rejuvenation of injection is that the ratio λ/r_g be small enough, and hence that the ratio of $\kappa_\perp/\kappa_\parallel$ be sufficiently large [but still less than unity: see Eq. (3)]. The critical value of λ/r_g is dependent on the downstream field angle Θ_{B2} to the shock normal in a non-trivial way. Physically meaningful values of λ/r_g can always be found to effect injection and efficient acceleration in highly oblique shocks^{5,20}. The survey work of Ellison et al.²⁰ found that for oblique shocks the efficiency of non-thermal ion generation declined rapidly with increasing Mach numbers and increasing values of λ/r_g , and further that unless the turbulence forced the situation very close (i.e. $\lambda/r_g \lesssim 2$) to the Bohm diffusion limit, there was a significant reduction in acceleration efficiency with increasing obliquity. Jokipii²⁹ observed that large values of λ/r_g reduced the timescale for acceleration in quasi-perpendicular shocks, and identified this property as a means to ease problems⁸ with accelerating cosmic rays all the way up to the “knee” using supernova remnants. Ellison et al.²⁰ highlighted the trade-off that then arises in oblique shocks: those that accelerate quickly are inefficient, and those shocks that produce cosmic rays copiously take longer to do so. This has a significant impact on models of cosmic ray production.

The appropriateness of small values of λ/r_g for highly oblique interplanetary shocks was affirmed by Baring et al.⁵, who obtained impressive fits to Ulysses distribution data for two shocks less than about 3 AU from the sun, using a Monte Carlo approach^{18,20}. Distributions were derived using upstream solar wind properties as input for the simulation. Such accurate fits, obtained for ion speeds up to around 4–5 times the solar wind speed, were possible only for values $\lambda/r_g \lesssim 3$ in shocks with $\Theta_{B1} \gtrsim 60^\circ$. Such low values are comparable to those deduced from spacecraft observations of interplanetary magnetic field turbulence^{26,27}. Kang and Jones³⁰ obtained similar fits in regimes of strong turbulence using an alternative theoretical approach, the numerical solution of the convection-diffusion differential equation for particle transport in shocks. An elegant aspect of the work of Baring et al.⁵ was the ability to simultaneously model proton and alpha-particle data; from this success, there is a strong suggestion that the particles interact with the field turbulence in a more-or-less elastic manner. Such theory/data comparisons are yet to be fully extended to shocks a little more distant from the sun; these possess the added complication of significant pick-up ion components^{5,25}. Note that interplanetary shocks have low Mach numbers, and are therefore weak; they are thus comparatively inefficient accelerators, so that using test-particle theories to model them is quite fitting.

The discussion so far has focused on non-relativistic shocks, for which there is an abundance of observational data in the heliosphere and applications in astrophysics. It is appropriate to comment briefly on results for relativistic shocks (where $u_1 \gtrsim 0.1c$), which are less comprehensively studied, yet may be quite important for active galaxies. Relativistic gases have adiabatic indices of $\gamma = 4/3$, and hence are more compressible than their non-relativistic counterparts: they can generate compression ratios as large as $r = 7$ in linear shock applications. Shocks with $u_1 \gtrsim 0.1c$ produce power-law distributions of relativistic particles, yet Eq. (2) cannot be used to predict their power-law index primarily because the assumption of isotropy must be relinquished. Relativistic shock fronts move sufficiently fast that even ultrarelativistic particles have difficulty catching them. Hence escape from the shock is greater and particle isotropy is impossible to achieve. At the same time, the energy boosts that ions receive in diffusion back and forth across the shock are enhanced relative to non-relativistic shock cases. The net effect is that for $f(p) \propto p^{-\sigma}$, the spectral index σ decreases with increasing shock speed u_1 . This behaviour was identified in the pitch-angle diffusion (i.e. imposing incremental changes in particle direction) transport equation analysis of Kirk and Schneider³¹ and the large angle scattering kinematic Monte Carlo simulation of Ellison et al.³². A comparison of their results³² reveals that the type of diffusion, i.e. pitch angle or large angle, strongly affects the population anisotropy and therefore is crucial to the determination of the power-law index.

The absence of the mention of electrons in the above discussions is striking. This is largely because the injection of electrons into the Fermi process is at present poorly understood. While protons at all energies from thermal upwards can resonate with Alfvén waves, it is not clear what waves thermal electrons can interact with to effect diffusion. Above around 20 MeV, electrons resonate with Alfvén waves and can actively participate in the diffusive acceleration process producing results³³ similar to those addressed above. However, between around 20 MeV and 0.5 MeV electrons can “scatter” off *whistler* waves (e.g. Levinson³⁴). Yet below this energy the picture is particularly unclear, and is complicated by the fact that whistlers can be strongly damped in hot plasmas; Alfvén modes usually escape this fate. Given ubiquitous observations of non-thermal radiation in astrophysics, as most non-thermal electronic radiative processes are so much more efficient than ionic ones, the inference of non-thermal (and predominantly relativistic) electron populations in numerous sources seems robust. Hence, if shock acceleration is assumed to be the origin such electrons, it remains to solve the problem of their injection into the Fermi mechanism. The role of electrons is discussed in the accompanying paper (Baring, this volume) that reviews applications of shock acceleration to gamma-ray production in young supernova remnants.

4 Non-Linear Modifications: a Brief Look

Linear test-particle models of shock acceleration are valid provided that the accelerated population does not alter the flow hydrodynamics. For strong, non-relativistic shocks $r = 4$, Eq. (2) indicates a p^{-2} distribution, which provides a divergent cosmic ray pressure or energy density if the momenta are permitted to extend to infinity. Consequently, if the maximum momentum achieved in a particular shock environment is large enough, the accelerated population dominates the gas pressure and can therefore modify or dominate the flow dynamics. This maximum p is generally determined by obvious natural limits such as the diffusion length of particles (greater than their gyroradii) being less than the physical size of the astrophysical system, beyond which particle escape ensues. The standard approach for describing such non-linear modifications to the flow is using the *Rankine-Hugoniot* relations for conservation of particle number, momentum, and energy fluxes. These relations form the central theme of the earliest approach to non-linear shock acceleration theory, the *two-fluid* (i.e. thermal gas plus cosmic rays) model³⁵, reviewed comprehensively by Drury¹⁶. The energy flux is the highest order momentum moment of the distribution: hence necessity of a non-linear treatment is contingent upon the pressure of the cosmic rays being a sizeable fraction of that of the thermal gas. Generally, the consideration of non-linear effects is imperative³⁶ for shocks with $r \gtrsim 3.4$.

The non-linearity is manifested through the feedback of the particles on the spatial profile of the flow velocity¹⁶, which in turn determines the shape of the distribution. The accelerated population presses on the upstream plasma and slows it down before the discontinuity is reached. An upstream *precursor* forms, in which the flow speed is monotonically decreasing. At the same time, the cosmic rays press on the downstream gas, slowing it down also. The net effect that usually emerges is one where the overall compression ratio, from far upstream to far downstream of the discontinuity, actually **exceeds that of the test-particle scenario**. This phenomenon was identified by Eichler³⁷, and Ellison and Eichler³⁶, and arises because of the possibility of particle escape at high energies; such a drain of energy and momentum flux must be compensated by ramping up the overall compression ratio to conserve the fluxes. It is illustrated and discussed in more detail in the accompanying paper by Baring in this volume (see also Jones and Ellison¹⁷). The highest energy particles generally have longer diffusion lengths³⁸ (unless κ is chosen to be independent of energy³⁹, an improbable situation), so that they sample a broader portion of the flow velocity profile, and hence larger compression ratios. Consequently, these particles have a flatter power-law index than those at lower energies, thereby dominating the pressure in a non-linear fashion. The resulting upward spectral curvature^{36,37} is the trademark of non-linear acceleration. Such spectral properties can be probed by Monte Carlo and plasma simulations, but cannot easily be explored using the two-fluid technique.

It remains to point out that such non-linear considerations are essential in certain real applications, even though they are seldom implemented in astrophysical models. The seminal work in this regard was performed by Ellison, Möbius and Paschmann³, who used the Monte Carlo technique to fit AMPTE data just upstream and downstream of quasi-parallel portions of the Earth's bow shock. The fit was impressive for thermal and non-thermal protons, and non-thermal He^{++} and a mixture of C,N and O ions (for which no thermal distributions could be measured). Furthermore, it *required* the implementation of non-linear modifications in this strong shock, and test-particle models were excluded by the data. This comparison has since been reproduced by other techniques, namely using hybrid plasma simulations^{21,40} and solutions to the convection-diffusion transport equation⁴¹. At the same time, the importance of non-linear considerations for electron acceleration have been outlined by Reynolds and Ellison³³, including the suggestion of an inference of spectral curvature from radio data for Kepler's supernova remnant. The role of non-linear shock modifications on electron acceleration is examined at greater length in accompanying paper (Baring, this volume) on SNR applications. The necessity

of non-linear considerations to supernova remnant modelling is underlined in the discussions there, augmenting the brevity of this exposition. To conclude, *non-linear* shock acceleration theory is central to the modelling of the Earth's bow shock and SNRs, and may also be extremely relevant to relativistic shock applications, territory that is yet to be explored.

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